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CONSTRAINTS ON COMPOSITE-MODELS EFFECTIVE LAGRANGIANS FROM $o\nu\beta\beta$ DECAY. *†

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ABSTRACT

We give a brief review of the existing bounds on effective couplings for excited states (e^*, ν^*, u^*, d^*) of the ordinary quarks and leptons arising from a composite model scenario. We then explore the phenomenological implications of the hypothesis that the excited neutrino ν^* might be of Majorana type. Recent bounds on the half-life of the $\Delta L = 2$ neutrinoless double beta decay $(0\nu\beta\beta)$ are used to constraint the compositeness effective couplings. We show that the bounds so obtained are roughly of the same order of magnitude as those available from high energy experiments.

1. Introduction

Up to now the behaviour of quarks and leptons has been, so far, very successfully described by a theory of pointlike quantum fields interacting through $SU(2) \times U(1) \times SU(3)_c$ gauge interactions, which is nowadays referred to as the standard model. We may however speculate that, when exploring higher energy ranges (from 1 TeV up to 15 TeV as planned with the next generation of supercolliders like LHC or NLC), we might hit an energy scale Λ_C at which a sub-structure of those "elementary" particles will show up, (see Fig. 1 for a schematic illustration). Although so far there is no experimental evidence signaling a further level of sub-structure, we cannot a priori exclude this possibility, and must explore its phenomenological consequences.

The idea that quarks and leptons might not be genuine elementary particles has been around for quite some time. Many models based on the idea that quarks and leptons are bound states of still unknown entities (generally referred to as *preons*), have already been proposed[1] but no completely consistent dynamical theory has been found

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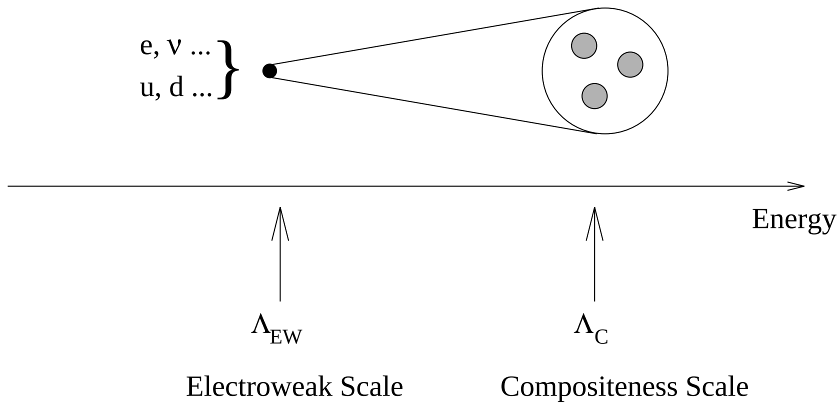

Fig. 1. The idea of compositeness: at the energy scale Λ_C (the compositeness scale, still unknown), ordinary quarks and leptons might show an internal structure.

so far. It is therefore particularly important to study model independent features of the idea of compositeness. There are two main consequences of having composite quarks and leptons:

(i) four-fermion contact interactions of dimension six;

(ii) highly massive excited states which couple to the ordinary fermions through gauge interactions;

Both facts are expected to give observable deviations from the predictions of the standard model provided that the compositeness scale Λ_C is not too large.

In this work we will first review what bounds on the compositeness scale can be derived from the study of the above mentioned effective interactions, and then we will show how current lower bounds for the half-life of the neutrinoless double beta decay ($0\nu\beta\beta$) can be used to get constraints on the compositeness effective couplings, when assuming the existence of a heavy composite Majorana neutrino.

2. Contact Interactions

In preon models, modifications to the gauge boson propagators and the interaction vertices with fermions are expected. To describe the former one can use[2] a simple parametrization by multiplying gauge boson propagators by a form factor $F(Q^2) \approx 1 + Q^2/\Lambda_C^2$. Such form factors effects were experimentally searched for, already in 1981-82 at PETRA, looking for deviations from the standard model predictions in the cross sections for $e^+e^- \to f\bar{f}$ with $(f = e, \mu, \tau, q)$. These experiments [3] gave lower bounds on Λ_C of the order of 100 GeV.

Composite fermions are also expected to have additional effective four fermions interactions through constituent exchange. Eichten, Lane and Peskin[4] proposed the following effective lagrangian to parametrize flavour diagonal helicity conserving contact interactions:

$$\mathcal{L}^{(ff)}_{Cont.} = \frac{g^2}{2\Lambda_C^2}\Big[\sum_{i,j=R,L} \eta_{ij}\, \bar{f}_j \gamma_\mu f_i\, \bar{f}_j \gamma^\mu f_j\Big] \tag{1}$$

where the compositeness scale Λ_C is defined in such a way that $g^2/4\pi = 1$ (i.e. the coupling g is strong) and $\max|\eta_{ij}| = 1$ and $i, j = L, R$. In Fig. 2 we show how contact interactions can affect fermion scattering at high energies. Clearly the interference between the contact term and the standard model diagram will give contributions that, relative to the standard model one, will be of the order:

$$\simeq (\frac{\alpha_i}{Q^2})^{-2} \frac{\alpha_i g^2}{Q^2 \Lambda_C^2} = \frac{g^2}{\alpha_i} \frac{Q^2}{\Lambda_C} \tag{2}$$

and will thus overwhelm the form factor contributions which are of the order Q^2/Λ_C^2.

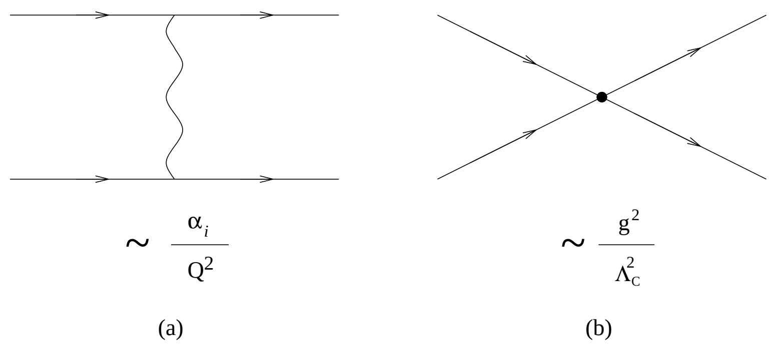

Fig. 2. Contact interactions give modifications to the fermion fermion scattering amplitude. The interference between the standard model diagram (a) and the one coming from contact interactions (b) will be the most important correction relative to the standard model prediction.

The effective lagrangian in Eq. (1) can be extended to include (model dependent) possibile non-diagonal contact interactions (which are expected if different fermions share common constituents). We could have:

$$\mathcal{L}^{(e\mu)}_{cont.} = \frac{g^2}{\Lambda_C^2} \sum_{i,j=R,L} \eta_{ij} \bar{e}_i \gamma_\lambda e_i \bar{\mu}_j \gamma^\lambda \mu_j \tag{3}$$

and similarly for (e, q).

Let us now discuss what bounds on the compositeness scale can be derived from the non-observation of deviations from the standard model predictions. We will use the following notation for the compositeness scale: $\Lambda^{\pm}_{ij}$ corresponds to the choice $\eta_{ij} = \pm 1$ and $\eta_{kl} = 0$ for $k \neq i, l \neq j$. High-energy fermion scattering, at electron positron and/or hadron colliders, has been used to study possible manifestations of the effective lagrangians given in Eqs. (1,3). Recent bounds come from the LEP experiment[5] at CERN where the study of the process $e^+e^- \to e^+e^-(\mu^+\mu^-)$ at the energy of the Z^0

Table 1. Lower bounds on the compositeness scale Λ_C from high energy fermion scattering and from leptonic tau decays.

process	Λ^+_{LL}(TeV) $>$	Λ^-_{LL}(TeV) $>$
$e^+e^- \to e^+e^-$	1.6	3.6
$e^+e^- \to \mu^+\mu^-$	2.6	1.9
$p\bar{p} \to e^+e^- + X$	1.7	2.2
$\tau \to \nu_\tau e \bar{\nu}_e$	3.8	8.1

resonance has given lower bounds on the compositeness scale: $\Lambda^+_{LL}(eeee) > 1.6$ TeV and $\Lambda^+_{LL}(ee\mu\mu) > 2.6$ TeV.

The Drell-Yan process ($p\bar{p} \to e^+e^- + X$) has been used at FERMILAB to obtain lower bounds on the compositeness scale of contact interactions between quarks and leptons[6]: $\Lambda^+_{LL}(eeqq) > 1.7$ TeV. Recently Diaz-Cruz and Sampayo[7] derived bounds on $\Lambda(\tau\nu_\tau e\nu_e)$ from a theoretical analysis of the effect of contact interactions in the leptonic τ decays, as shown in Fig. 3.

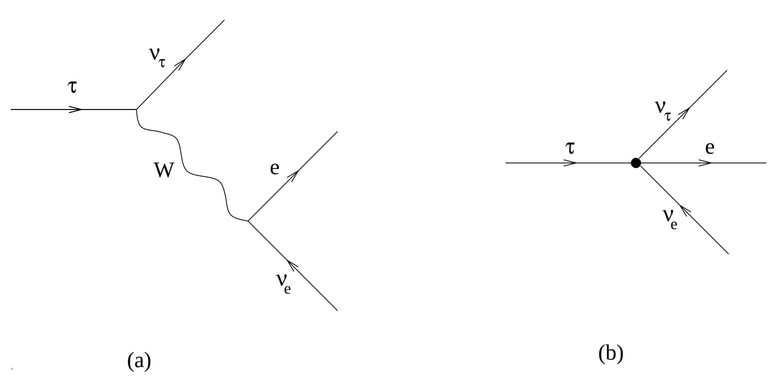

Fig. 3. Leptonic τ decays: (a) standard model contribution; (b) contact interactions contributuion.

They find a rather interesting bound; $\Lambda^+_{LL}(\tau\nu_\tau e\nu_e) > 3.8$ TeV but it is to be remarked that their result is based on the assumption that the compositeness scale is flavour dependent, which leads to $\Lambda(\tau\nu_\tau e\nu_e) \ll \Lambda(\mu\nu_\mu e\nu_e)$. The above mentioned bounds are summarized in Table 1. We also refer the reader to the Review of Particle

Properties of the Particle Data Group [8] for a list of previous bounds.

3. Excited states of ordinary fermions

Although no completely consistent dynamical composite theory has been found up to now, one inevitable common prediction of composite models is the existence of excited states of the known quarks and leptons, analogous to the series of higher energy levels of the hydrogen atom. The masses of the excited particles should not be much lower than the compositeness scale Λ_C, which is expected to be at least of the order of a TeV according to the experimental constraints discussed in the previous section. We expect therefore the heavy fermion masses to be, *at least*, of the order of a few hundred GeV.

Excited leptons (l^*) and quarks (q^*) are expected to interact with light fermions via gauge interactions. Let us consider the coupling $e^*e\gamma$, depicted in Fig. 4 and for which Low[9] proposed in 1965 a magnetic moment type interaction:

$$\mathcal{L}_{int} = \frac{ef_\gamma}{2m^*}\bar{\psi}_{e^*}\sigma_{\mu\nu}\psi_e F^{\mu\nu} + \text{h.c.} \tag{4}$$

where m^* is the mass of the excited electron and f_γ is a dimensionless coupling constant; $F_{\mu\nu}$ is the electromagnetic field strength tensor.

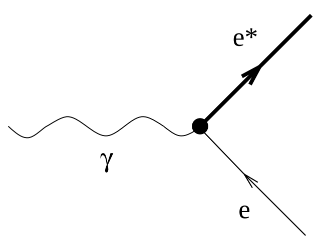

Fig. 4. The transition coupling between a light electron to its correspondent excited state (e^*) via gauge interactions.

Clearly, this coupling can produce deviations from the predictions of the standard model through the exchange of virtual heavy excited states, or it may be responsible for the direct production and decay of the excited states in high energy fermion scattering. These effects can be used to derive bounds on the coupling constants appearing in Eq. 4. In 1982 Renard [10] showed that the precise measurements of the anomalous electron magnetic moment give bounds on the masses of the excited states (or equivalently the compositeness scale). In Fig. 5 we show one of the diagrams involving the exchange of a virtual excited electron and contributing to the electron's anomalous magnetic moment and electric dipole moment. Renard considered a general tensor and pseudo-tensor effective coupling:

$$\mathcal{L}_{int} = \frac{e}{2m^*}\bar{\psi}_{e^*}\,\sigma_{\mu\nu}(a - ib\gamma_5)\psi_e\,F^{\mu\nu} + \text{h.c.} \tag{5}$$

which is a simple generalization of Eq. 4.

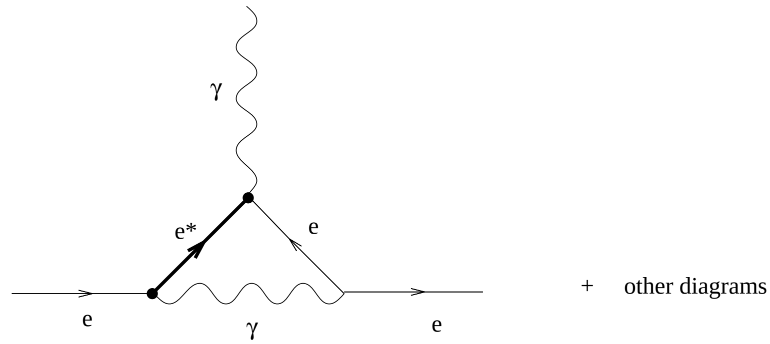

Fig. 5. One of the diagrams contributing to the electron anomalous magnetic moment, in addition to those of the standard model, in the hypothesis that excided states exist.

Defining the anomalous magnetic moment and electric dipole moment of the electron through the lagrangian:

$$\mathcal{L} = e\bar{\psi}_e \Big\{ \gamma^\mu A_\mu + \frac{\chi}{4m_e}\sigma^{\mu\nu}F_{\mu\nu} - i\frac{\chi'}{4m_e}\sigma^{\mu\nu}F_{\mu\nu}\gamma^5 \Big\} \bar{\psi}_e \tag{6}$$

in the limit $m_{e^*} \gg m_e$ one obtains [10]:

$$\begin{aligned} \chi &= \frac{4\alpha}{\pi}(|a|^2 - |b|^2)\frac{m_e}{m_{e^*}} + \frac{9\alpha}{2\pi}(|a|^2 + |b|^2)\frac{m_e^2}{m_{e^*}^2}\,; \\ \chi' &= \frac{8\alpha}{\pi}\Re(ab^*)\frac{m_e}{m_{e^*}}\,. \end{aligned} \tag{7}$$

The small value of the electron dipole moment $d_e = \chi'/2m_e \simeq 0.7 \times 10^{-24}$ cm implies that a and b in Eq. 5 cannot be simultaneously real: $\Re(ab^*) \simeq 0$. Regarding the anomalous magnetic moment contribution, we see that the effect appears at order m_e/m_{e^*} if $|a| \neq |b|$, while if there is *chiral symmetry* ($|a| = |b|$) the effect is of order $(m_e/m_{e^*})^2$. Thus, the precise measurement of the electron's $g-2$ will put a weaker or stronger contraint on the value of m_{e^*} depending on whether or not the coupling respects chiral symmetry. From $\delta\chi_e \leq 2 \times 10^{-10}$ one finds:

$$\begin{aligned} m_{e^*} &\geq (|a|^2 - |b|^2) \times 22\,\mathrm{TeV} && \text{without chiral symmetry} \\ m_{e^*} &\geq (|a|^2 + |b|^2)^{1/2} \times 3.8\,\mathrm{GeV} && \text{with chiral symmetry} \end{aligned} \tag{8}$$

the corresponding bounds for the muon are:

$$\begin{aligned} m_{\mu^*} &\geq (|a|^2 - |b|^2) \times 110\,\mathrm{TeV} && \text{without chiral symmetry} \\ m_{\mu^*} &\geq (|a|^2 + |b|^2)^{1/2} \times 110\,\mathrm{GeV} && \text{with chiral symmetry} \end{aligned} \tag{9}$$

One is thus led to conclude that if the compositeness scale Λ_C ($\approx m^*$) is of the order of one to a few TeV, then the coupling in Eq. 5 must display chiral symmetry, i.e.

the heavy excited state can couple only to a left-handed (or only to a right-handed) light fermion.

The extension of the effective coupling given in Eq. 5 including electroweak interactions has been discussed in the literature [11] using weak isospin (I_W) and hypercharge (Y) conservation. Within this model, it is assumed that the lightness of the ordinary leptons could be related to some global unbroken chiral symmetry which would produce massless bound states of preons in the absence of weak perturbations due to $SU(2) \times U(1)$ gauge and Higgs interactions. The large mass of the excited leptons arises from the unknown underlying dynamics and *not* from the Higgs mechanism.

Assuming that such states are grouped in $SU(2) \times U(1)$ multiplets, since light fermions have $I_W = 0, 1/2$ and electroweak gauge bosons have $I_W = 0, 1$, only multiplets with $I_W \leq 3/2$ can be excited in the lowest order in perturbation theory. Also, since none of the gauge fields carry hypercharge, a given excited multiplet can couple only to a light multiplet with the same Y.

In addition, conservation of the electromagnetic current forces the transition coupling of heavy-to-light fermions to be of the magnetic moment type respect to any electroweak gauge bosons [11]. In fact, a γ_μ transition coupling between e and e^* mediated by the $\vec{W}^\mu$ and B^μ gauge fields, would result in an electromagnetic current of the type $j^\mu_{e.m.} \approx \bar{\psi}_{e^*}\gamma^\mu\psi_e$ wich would not be conserved due to the different masses of excited and ordinary fermions, (actually it is expected that $m_{e^*} \gg m_e$).

Let us here restrict to the first family and consider spin-1/2 excited states grouped in multiplets with $I_W = 1/2$ and $Y = -1$,

$$L^* = \begin{pmatrix} \nu^* \\ e^* \end{pmatrix} \tag{10}$$

which can couple to the light left-handed multiplet

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} = \frac{1-\gamma_5}{2}\begin{pmatrix} \nu \\ e \end{pmatrix} \tag{11}$$

through the gauge fields $\vec{W}^\mu$ andB^μ. The relevant interaction can be written[11] in terms of two *new* independent coupling constants f and f':

$$\begin{aligned} \mathcal{L}_{int} \;\; = \;\; & \frac{gf}{\Lambda_C}\bar{L}^*\sigma_{\mu\nu}\frac{\vec{\tau}}{2}l_L \cdot \partial^\nu \vec{W}^\mu \\ & \qquad + \frac{g'f'}{\Lambda_C}\left(-\frac{1}{2}\bar{L}^*\sigma_{\mu\nu}l_L\right) \cdot \partial^\nu B^\mu + \text{h.c.} \end{aligned} \tag{12}$$

where $\vec{\tau}$ are the Pauli $SU(2)$ matrices, g and g' are the usual $SU(2)$ and $U(1)$ gauge coupling constants, and the factor of $-1/2$ in the second term is the hypercharge of the $U(1)$ current. This effective lagrangian has been widely used in the literature to predict production cross sections and decay rates of the excited particles [8,12,13].

In terms of the physical gauge fields $Z_\mu = -B_\mu \sin\theta_W \; + W^3_\mu \cos\theta_W$, $A_\mu = B_\mu \cos\theta_W + W^3_\mu \sin\theta_W$ and $W^\pm_\mu = (1/\sqrt{2})\big(W^1_\mu \mp i\, W^2_\mu\big)$, the effective interaction in Eq. 12 can be reexpressed as

$$\mathcal{L}_{int} = \sum_{V=\gamma,Z,W} \frac{e}{\Lambda_{\rm C}} \bar{f}^* \, \sigma_{\mu\nu}(c_{Vf^*f} - d_{Vf^*f}\gamma_5)\, f\, \partial_\mu V_\nu + \text{h.c.} \tag{13}$$

where the coupling constants have to satisfy the condition $|c_{Vf^*f}| = |d_{Vf^*f}|$, if we require chiral symmetry, and are related to f and f' by the following relations:

$$\begin{aligned}
c_{\gamma f^* f} &= -\frac{1}{4}(f + f') \\
c_{Z f^* f} &= -\frac{1}{4}(f \cot\theta_W + f' \tan\theta_W) \\
c_{W e^* \nu} &= \frac{f}{2\sqrt{2}\sin\theta_W} \\
c_{\gamma \nu^* \nu} &= -\frac{1}{4}(f - f') \\
c_{Z \nu^* \nu} &= -\frac{1}{4}(f \cot\theta_W - f' \tan\theta_W) \\
c_{W \nu^* e} &= \frac{f}{2\sqrt{2}\sin\theta_W}
\end{aligned} \tag{14}$$

The extension to quarks and strong interactions as well as to other multiplets and a detailed discussion of the spectroscopy of the excited particles can be found in the literature [14].

Here, let us we write down explicitly the interaction lagrangian describing the coupling of the heavy excited neutrino with the light electron, as it will be used in the following section in order to discuss bounds on the compositeness effective couplings from low-energy, nuclear, double-beta decay:

$$\mathcal{L}_{eff} = (\frac{gf}{\sqrt{2}\Lambda_{\rm C}})\Big\{\Big(\overline{\nu}^* \sigma^{\mu\nu} \frac{1-\gamma_5}{2}\, e\Big)\, \partial_\nu W^+_\mu\Big\}\,. \tag{15}$$

Let us now discuss the current bounds on the mass of the excited states derived from fermion scattering experiments. Limits on m_{e^*} can be obtained from indirect effects due to t-channel e^* exchange in $e^+e^- \to \gamma\gamma$. The L3 collaboration has found [15] at $\sqrt{s} = 91$ GeV (LEP):

$$m_{e^*} > 127 \text{ GeV}. \tag{16}$$

Limits on the masses of the excited leptons $(e^*\,, \nu^*)$ from single production in electron-positron collisions $e^+e^- \to e^* e$, $e^+e^- \to \nu^* \nu$ at a center of mass energy corresponding to the Z^0 resonance, are roughly given by the Z^0 mass[16]:

$$m_{e^*,\nu^*} > 91\,\text{GeV} \tag{17}$$

Pair-production of l^*, q^* rely on the electroweak charge of the excited particles and gives usually less constraining [8] bounds. These constraints are obviously limited by the center-of-mass energy of the accelerator.

As regards the bounds on the masses of the excited quarks, the strongest comes from an analysis of the reaction $p\bar{p} \to q^*X$ with $q^* \to q\gamma, qW$ at a center of mass energy of 1.8 TeV. Assuming u^* and d^* to be degenerate and $f, f', f_s = 1$ (where f_s is the dimensionless coupling constant of the strong transition magnetic coupling q^*qg [14], corresponding to f and f' appearing in Eq. 12) the CDF collaboration [17] has found:

$$m_{q^*} > 540\,\mathrm{GeV}. \tag{18}$$

The above bounds are summarized in Table 2 and we refer the reader to ref. [8] for a list of earlier bounds.

Table 2. Current lower bounds on the masses of the excited states of ordinary fermions as they are deduced from direct search high energy experiments.

excited particle	$m_*(\mathrm{GeV}) >$	process
e^*	127.	$e^+e^- \to \gamma\gamma$
e^*	91.	$e^+e^- \to e^+e^*$
ν^*	91.	$e^+e^- \to \nu^+\nu^*$
q^*	540.	$p\bar{p} \to q^* + X$

Finally, let us remark that the ZEUS and H1 collaborations (DESY) have recently published [18,19] the results of a search of excited states from single production in electron-proton collisions at HERA. They have studied the reaction $ep \to l^*X$ with the subsequent decay $l^* \to l'V$ where $V = \gamma, Z, W$ (see Fig. 6). Upper limits for the quantity $\sqrt{|c_{Vl^*e}|^2 + |d_{Vl^*e}|^2}/\Lambda_{\mathrm{C}} \times \mathrm{Br}^{1/2}(l^* \to l'V)$ are derived[18] as a function of the excited lepton mass and for the various decay channels. They are sensitive up to 180 GeV for m_{ν^*} and up to 250 GeV for m_{e^*,q^*}.

For the purpose of comparing our recent analysis of double-beta decay bounds on compositeness[20] with the bounds discussed above, (see section 4) we quote here the limit on the ν^* coupling that the ZEUS collaboration has obtained at the highest

accessible mass ($m_{\nu^*} = 180$ GeV):

$$\frac{\sqrt{|c_{W\nu^*e}|^2 + |d_{W\nu^*e}|^2}}{\Lambda_C} \times \mathrm{Br}^{1/2}(\nu^* \to \nu W) \leq 5 \times 10^{-2}\,\mathrm{GeV}^{-1}. \tag{19}$$

Let us note that these limits depend on the branching ratios of the decay channel chosen.

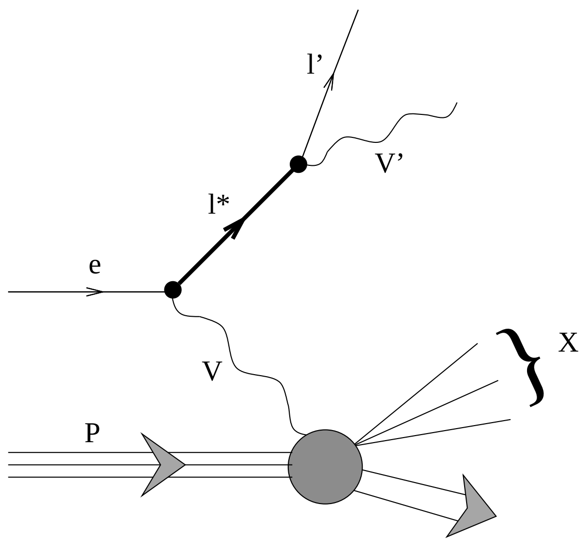

Fig. 6. Electroproduction of excited leptons $l^* = e^*, \nu^*$ through t-channel exchange of electroweak gauge bosons $V = \gamma, Z^0, W$.

4. Neutrinoless Double Beta Decay (o$\nu\beta\beta$).

In this section we discuss the possibility that the heavy excited neutrino ν^* (hereafter denoted N) might be a Majorana particle[20,21] and explore its low energy manifestations, namely neutrinoless double beta decay.

Heavy neutral Majorana particles with masses in the TeV region are predicted in various theoretical models, such as superstring-inspired E_6 grand unification [22] or left-right symmetric models [23]. In addition the possibility of a fourth generation with a heavy neutral lepton, that could be of Majorana type, is not yet ruled out [24,25].

In practical calculations of production cross sections and decay rates of excited states, it has been customary[13,26,27] to assume that the dimensionless couplings f and f' in Eq. (12) are of order unity. However if we assume that the excited neutrino is of Majorana type, we have to verify that this choice is compatible with present experimental limits on neutrinoless double beta decay ($0\nu\beta\beta$):

$$A(Z) \to A(Z+2) + e^- + e^- \tag{20}$$

a nuclear decay, that has attracted much attention both from particle and nuclear physicists because of its potential to expose lepton number violation. More generally, it is expected to give interesting insights about certain gauge theory parameters such as leptonic charged mixing matrix, neutrino masses etc. The process in Eq. (20),

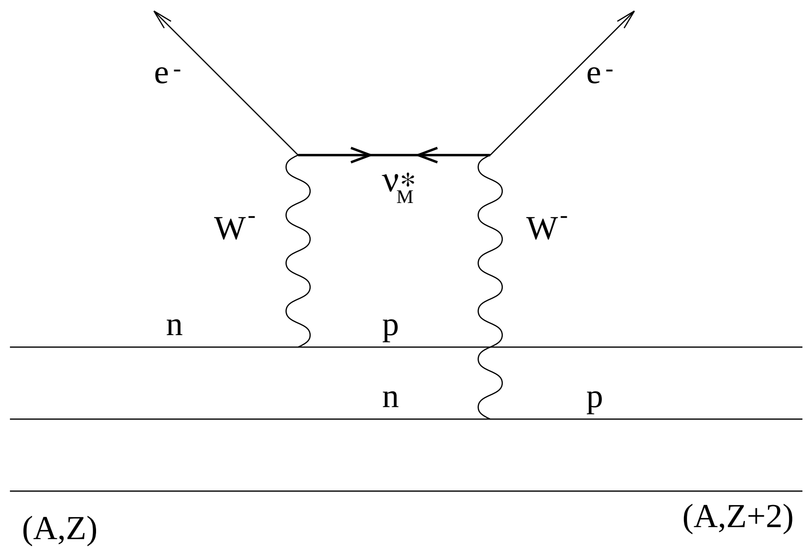

Fig. 7. Neutrinoless double beta decay ($\Delta L = +2$ process) mediated by a composite heavy Majorana neutrino.

which can only proceed via the exchange of a massive Majorana neutrino, has been experimentally searched for in a number of nuclear systems [28] and has also been extensively studied from the theoretical side [29,30,31].

We will consider here the decay

$$^{76}\mathrm{Ge} \to \, ^{76}\mathrm{Se} + 2e^- \tag{21}$$

for which we have from the Heidelberg-Moscow $\beta\beta$-experiment the recent limit [32] ($T_{1/2}$ is the half life = log2 × lifetime)

$$T_{1/2}\,(\,^{76}\mathrm{Ge} \to \, ^{76}\mathrm{Se} + 2e^-) \geq 5.1 \times 10^{24}\,\mathrm{yr} \quad 90\,\%\ \mathrm{C.L.} \tag{22}$$

In the following we estimate the constraint imposed by the above measurement on the coupling (f/Λ_C) of the heavy composite neutrino, as given by Eq. (15). The fact that neutrinoless double beta decay measurements might constrain composite models, was also discussed in ref.[21] but within the framework of a particular model and referring to a heavy Majorana neutrino with the usual γ_μ coupling. Models in which $0\nu\beta\beta$ decay proceeds via the exchange of a heavy sterile Majorana neutrino (mass in the GeV scale or higher) have also been recently considered [33].

The transition amplitude of $0\nu\beta\beta$ decay is calculated according to the interaction lagrangian:

$$\mathcal{L}_{int} \quad = \quad \frac{g}{2\sqrt{2}}\Big\{\frac{f}{\Lambda_C}\bar{\psi}_e(x)\sigma_{\mu\nu}(1+\gamma_5)\psi_N(x)\partial^\mu W^{\nu(-)}(x)$$

$$+\cos\theta_C J^h_\mu(x) W^{\mu(-)}(x) + h.c.\Big\} \tag{23}$$

where θ_C is the Cabibbo angle ($\cos\theta_C = 0.974$) and J^h_μ is the hadronic weak charged current.

We emphasize that in Eq. (23) we have a $\sigma_{\mu\nu}$ type of coupling as opposed to the γ_μ coupling so far encountered in all $0\nu\beta\beta$ decay calculations (see the discussion in section 3). For simplicity, we carry out our analysis assuming that there are no additional contributions to $0\nu\beta\beta$ decay from light Majorana neutrinos, right handed currents or other heavy Majorana neutrinos originating from another source.

The transition amplitude is then

$$\begin{aligned} S_{fi} &= (cos\theta_C)^2 (\frac{g}{2\sqrt{2}})^4 \left(\frac{f}{\Lambda_C}\right)^2 (\frac{1}{2}) \int \frac{d^4k}{(2\pi)^4}\, d^4x\, d^4y e^{-ik\cdot(x-y)} \times \\ & \frac{1}{\sqrt{2}}(1-P_{12})\bar{u}(p_1)\sigma_{\mu\lambda}(1+\gamma_5)\frac{\not k + M_N}{k^2 - M_N^2}(1+\gamma_5)\sigma_{\nu\rho} v(p_2) \times \\ & [F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2} e^{ip_1\cdot x} e^{ip_2\cdot y} \times \\ & (k-p_1)^\lambda (k+p_2)^\rho \frac{\langle f|J^\mu_h(x)\; J^\nu_h(y)|i\rangle}{[(k-p_1)^2 - M_W^2][(k+p_2)^2 - M_W^2]} \end{aligned} \tag{24}$$

where $(1-P_{12})/\sqrt{2}$ is the antisymmetrization operator due to the production of two identical fermions, the functions $F(Z,\epsilon)$ are the well known Fermi functions [34] that describe the distorsion of the electron's plane wave due to the nuclear Coulomb field (ϵ_i are the electron's kinetic energies in units of $m_e c^2$),

$$\begin{aligned} F(Z,\epsilon) &= \chi(Z,\epsilon)\frac{\epsilon+1}{[\epsilon(\epsilon+2)]^{1/2}} \\ \chi(Z,\epsilon) &\approx \chi^{R.P.}(Z) = \frac{2\pi\alpha Z}{1-e^{-2\pi\alpha Z}} \quad \text{(Rosen-Primakoff approximation)} \end{aligned} \tag{25}$$

As is standard in such calculations, we make the following approximations [29,30]:
i) the hadronic matrix element is evaluated within the closure approximation

$$\langle f|J^\mu_h(x)\; J^\nu_h(y)|i\rangle \approx e^{i(E_f - \langle E_n\rangle)x_0} e^{i(\langle E_n\rangle - E_i)y_0} \langle f|J^\mu_h(\boldsymbol{x})\; J^\nu_h(\boldsymbol{y})|i\rangle \tag{26}$$

where $\langle E_n\rangle$ is an average excitation energy of the intermediate states. This allows one to perform the integrations over k_0, x_0, y_0 in Eq. (24);
ii) neglect the external momenta p_1, p_2 in the propagators and use the long wavelength approximation: $\exp(-i\boldsymbol{p}_1\cdot\boldsymbol{x}) = \exp(-i\boldsymbol{p}_2\cdot\boldsymbol{x}) \approx 1$;
iii) the average virtual neutrino momentum $\langle|\boldsymbol{k}|\rangle \approx 1/R_0 = 40$ MeV is much larger than the typical low-lying excitation energies, so that, $k_0 = E_f + E_1 - \langle E_n\rangle$ can be neglected relative to $\boldsymbol{k}$;
iv) the effect of W and N propagators can be neglected since $M_W \approx 80$ GeV is much

greater than $\boldsymbol{k}$ in the region where the integrand is large, and we are interested in heavy neutrino masses $M_N \gg M_W$.

For the hadronic current we make the usual ansatz:

$$\begin{aligned} J^h_\mu(\boldsymbol{x}) &= \sum_k j_\mu(k)\delta^3(\boldsymbol{x}-\boldsymbol{r}_k) \\ j_\mu(k) &= \overline{\mathcal{N}}_k\gamma_\mu(f_V - f_A\gamma_5)\tau_+(k)\mathcal{N}_k\, f_A(|\boldsymbol{q}|^2) \end{aligned} \tag{27}$$

where $\mathbf{r}_k$ is the coordinate of the k-th nucleon, $\tau_+(k) = (1/2)(\tau_1(k) + i\tau_2(k))$ is the step up operator for the isotopic spin, ($\vec{\tau}(k)$ is the matrix describing the isotopic spin of the k-th nucleon), $\mathcal{N} = \binom{\psi_p}{\psi_n}$ and we have introduced the nucleon form factor:

$$f_A(|\boldsymbol{q}|^2) = \frac{1}{(1+|\boldsymbol{q}|^2/m_A^2)^2}, \tag{28}$$

with $m_A = 0.85$ GeV, in order to take into account the finite size of the nucleon, which is known to give important effects for the heavy neutrino case. We also take the nonrelativistic limit of the nuclear current:

$$j_\mu(k) = f_A(|\boldsymbol{q}|^2)\times\tilde{j}_\mu(k) \qquad \tilde{j}_\mu(k) = \begin{cases} f_V\tau_+(k) & \text{if } \mu=0 \\ -f_A\tau_+(k)(\sigma_k)_i & \text{if } \mu=i \end{cases} \tag{29}$$

($\vec{\sigma}_k$ is the spin matrix of the k-th nucleon). Then using the same notation as in Ref. [30] we arrive at

$$\begin{aligned} S_{fi} &= (G_F\cos\theta_C)^2\frac{f^2}{\Lambda_C^2}\frac{1}{2}2\pi\delta(E_0-E_1-E_2)\times \\ &\quad \frac{1}{\sqrt{2}}(1-P_{12})\bar{u}(p_1)\sigma_{\mu i}\sigma_{\nu j}(1+\gamma_5)v(p_2)[F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2}\times \\ &\quad M_N\sum_{kl} I_{ij}\,\langle f|\tilde{j}^\mu(k)\tilde{j}^\nu(l)|i\rangle \end{aligned} \tag{30}$$

where I_{ij} is an integral over the virtual neutrino momentum,
$(\boldsymbol{r}_{kl} = \boldsymbol{r}_k - \boldsymbol{r}_l\,, r_{kl} = |\boldsymbol{r}_k - \boldsymbol{r}_l|\,, x_{kl} = m_A r_{kl})$

$$\begin{aligned} I_{ij} &= \frac{1}{M_N^2}\int\frac{d^3\boldsymbol{k}}{(2\pi)^3}\frac{(-k_ik_j)}{(1+|\boldsymbol{k}|^2/m_A^2)^4}\exp{(i\boldsymbol{k}\cdot\boldsymbol{r}_{kl})} \\ &= \frac{1}{4\pi}\frac{m_A^4}{M_N^2}\frac{1}{r_{kl}}\left\{-\delta_{ij}F_A(x_{kl}) + \frac{(\boldsymbol{r}_k)_i(\boldsymbol{r}_l)_j}{r_{kl}^2}F_B(x_{kl})\right\} \end{aligned} \tag{31}$$

with:

$$\begin{aligned} F_A(x) &= \frac{1}{48}e^{-x}\,(x^2+x) \\ F_B(x) &= \frac{1}{48}e^{-x}\,x^3 \end{aligned} \tag{32}$$

Note that we can make the replacement
$\sigma_{\mu i}\sigma_{\nu j} \to (1/2)\{\sigma_{\mu i}, \sigma_{\nu j}\} = \eta_{\mu\nu}\eta_{ij} - \eta_{i\nu}\eta_{i\mu} + i\gamma_5\epsilon_{\mu i\nu j}$ since I_{ij} is a symmetric tensor and, with straightforward algebra, we obtain

$$\begin{aligned} S_{fi} &= M_{fi}\, 2\pi\, \delta(E_0 - E_1 - E_2) \\ M_{fi} &= (G_F \cos\theta_C)^2 \frac{1}{4}\frac{-1}{2\pi}\frac{f_A^2}{r_0 A^{1/3}}\, l\, \langle m \rangle \end{aligned} \tag{33}$$

where we have defined

$$\begin{aligned} l &= \frac{1}{\sqrt{2}}(1 - P_{12})\bar{u}(p_1)(1+\gamma_5)v(p_2)[F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2} \\ \langle m \rangle &= m_e \eta_N \langle f \,|\, \Omega \,|\, i \rangle \\ \eta_N &= \frac{m_p}{M_N} m_A^2 \left(\frac{f}{\Lambda_C}\right)^2 \\ \Omega &= \frac{m_A^2}{m_p m_e}\sum_{k\neq l} \tau_+(k)\tau_+(l)\frac{R_0}{r_{kl}}\left[\left(\frac{f_V^2}{f_A^2} - \vec{\sigma}_k\cdot\vec{\sigma}_l\right)(F_B(x_{kl}) - 3F_A(x_{kl}))\right. \\ &\qquad \left. -\vec{\sigma}_k\cdot\vec{\sigma}_l\, F_A(x_{kl}) + \frac{\vec{\sigma}_k\cdot \boldsymbol{r}_{kl}\,\vec{\sigma}_l\cdot\boldsymbol{r}_{kl}}{r_{kl}^2}F_B(x_{kl})\right] \end{aligned} \tag{34}$$

and $R_0 = r_0 A^{1/3}$ is the nuclear radius ($r_0 = 1.1$ fm).

The new result here is the nuclear operator Ω which is substantially different from those so far encountered in $0\nu\beta\beta$ decays, due to the $\sigma_{\mu\nu}$ coupling of the heavy neutrino that we are considering. The decay width is obtained upon integration over the density of final states of the two-electron system

$$d\Gamma = \sum_{final\, spins} |M_{fi}|^2\, 2\pi\delta(E_0 - E_1 - E_2)\frac{d^3\boldsymbol{p}_1}{(2\pi)^3\, 2E_1}\frac{d^3\boldsymbol{p}_2}{(2\pi)^3\, 2E_2} \tag{35}$$

and the total decay rate Γ can be cast in the form

$$\Gamma = (G_F\cos\theta_C)^4 \frac{(f_A)^4\, m_e^7\, |\eta_N|^2}{(2\pi)^5 r_0^2 A^{2/3}}\, f_{0\nu}(\epsilon_0, Z)\, |\Omega_{fi}|^2 \tag{36}$$

$$f_{0\nu} = \xi_{0\nu} f_{0\nu}^{R.P.} \tag{37}$$

$$f_{0\nu}^{R.P.} = |\chi^{R.P.}(Z+2)|^2 \frac{\epsilon_0}{30}(\epsilon_0^4 + 10\epsilon_0^3 + 40\epsilon_0^2 + 60\epsilon_0 + 30) \tag{38}$$

where, $\Omega_{fi} = \langle f|\Omega|i\rangle$, ϵ_0 is the kinetic energy of the two electrons in units of $m_e c^2$, and $\xi_{0\nu}$ is a numerical factor that corrects for the Rosen-Primakoff approximation [30] used in deriving the analytical expression of $f_{0\nu}^{R.P.}$. For the decay considered in Eq.(21), we have [30] $\xi_{0\nu} = 1.7$ and $\epsilon_0 = 4$. The half-life is finally written as

$$\begin{aligned} T_{1/2} &= \frac{K_{0\nu}\, A^{2/3}}{f_{0\nu}\, |\eta_N|^2\, |\Omega_{fi}|^2} \\ K_{0\nu} &= (\log 2)\frac{(2\pi)^5}{(G_F\cos\theta_C m_e^2)^4}\frac{(m_e r_0)^2}{m_e f_A^4} = 1.24\times\, 10^{16}\,\mathrm{yr} \end{aligned} \tag{39}$$

Combining Eq. (39) with the experimental limit given in Eq. (22), we obtain a constraint on the quantity $|f|/(\Lambda_{\rm C}^2 M_N)^{1/2}$

$$\frac{|f|}{(\Lambda_{\rm C}^2 M_N)^{1/2}} < \left(\frac{1}{m_p m_A^2}\right)^{1/2} \left[\frac{K_{0\nu}\, A^{2/3}}{5.1\times 10^{24}\,{\rm yr}\times f_{0\nu}(Z,\epsilon_0)}\right]^{1/4} \frac{1}{|\Omega_{fi}|^{1/2}} \tag{40}$$

Given the heavy neutrino mass M_N and the compositeness scale $\Lambda_{\rm C}$, we only need to evaluate the nuclear matrix element Ω_{fi} to know the upper bound on the value of $|f|$ imposed by neutrinoless double beta decay.

The evaluation of the nuclear matrix elements was in the past regarded as the principal source of uncertainty in $0\nu\beta\beta$ decay calculations, but the recent high-statistics measurement [35] of the allowed $2\nu\beta\beta$ decay, a second order weak-interaction β decay, has shown that nuclear physics can provide a very good description of these phenomena, giving high reliability to the constraints imposed by $0\nu\beta\beta$ decay on non-standard model parameters.

Since we simply want to estimate the order of magnitude of the constraint in Eq. (40) we will evaluate the nuclear matrix element only approximatly. First of all the expression of the nuclear operator in Eq. (34) is simplified making the following replacement [36]

$$\frac{r_{kl}^i r_{kl}^j}{r_{kl}^2} \rightarrow \langle \frac{r_{kl}^i r_{kl}^j}{r_{kl}^2} \rangle \rightarrow \frac{1}{3}\delta_{ij} \tag{41}$$

The operator Ω becomes then

$$\Omega \approx \frac{m_A^2}{m_p m_e}(m_A R_0)\sum_{k\neq l}\tau_+(k)\tau_+(l)\left(\frac{f_V^2}{f_A^2} - \frac{2}{3}\vec{\sigma}_k\cdot\vec{\sigma}_l\right)F_N(x_{kl}) \tag{42}$$

where $F_N = (1/x)(F_B - 3F_A) = (1/48)e^{-x}(x^2 - 3x - 3)$ with F_B and F_A given in Eq.(31).

Since we are interested in deriving the lowest possible upper bound on $|f|$ given by Eq. (40), let us find the maximum absolute value of the nuclear matrix element of the operator Ω in Eq.(42):

$$|\Omega_{fi}| \leq \frac{m_A^2}{m_p m_e}(m_A R_0)|F_N(\bar{x})|\left\{\frac{f_V^2}{f_A^2}|M_F| + \frac{2}{3}|M_{GT}|\right\} \tag{43}$$

where $M_F = \langle f|\sum_{k\neq l}\tau_+(k)\tau_+(l)|i\rangle$ and $M_{GT} = \langle f|\sum_{k\neq l}\tau_+(k)\tau_+(l)\vec{\sigma}_k\cdot\vec{\sigma}_l|i\rangle$ are respectively the matrix elements of the Fermi and Gamow-Teller operators whose numerical values for the nuclear system under consideration are [29,30], $M_F = 0$ and $M_{GT} = -2.56$. Inspection of the radial function F_N (for $x \geq 0$) shows that its maximum absolute value is attained at $x = 0$. In Eq. (43) we have evaluated F_N at $x = 2.28\,(r_{kl} = 0.5$ fm). This value of r_{kl} corresponds to the typical internuclear distance at which short range nuclear correlations become important [29], so that the

Table 3. Most stringent, lower bounds on Λ_C with $|f| = 1$, and upper bounds on $|f|$ with $\Lambda_C = 1$ TeV, for different values of the heavy neutrino mass M_N, as can be derived from the $0\nu\beta\beta$ half-life lower limit in Eq. (6), within the approximation discussed in the text.

M_N(TeV)		0.6	0.8	1.0	1.2	1.4	1.6	1.8	2.0
Λ_C (TeV) $>$	$[\lvert f\rvert = 1]$	0.33	0.29	0.26	0.23	0.22	0.20	0.19	0.18
$\lvert f\rvert <$	$[\Lambda_C = 1\,\text{TeV}]$	3.0	3.5	3.9	4.2	4.6	4.9	5.2	5.5

region $x \leq 2.28$ does not give contributions to the matrix element of the nuclear operator. We thus find:

$$|\Omega_{fi}| \leq 0.6 \times 10^3, \tag{44}$$

which together with Eq. (40) gives:

$$\frac{|f|}{\Lambda_C (M_N)^{1/2}} \leq 3.9 \text{ TeV}^{-3/2}. \tag{45}$$

However, since we have used an upper bound for the nuclear matrix element (Eq. (44)), the above should be taken as the most stringent upper bound one could possibly get for the quantity $|f|/(\Lambda_C^2 M_N)^{1/2}$ given the half-life measurement quoted in Eq. (22). An exact evaluation of the nuclear matrix element will give less stringent constraints than those that can be derived from Eq. (45).

With this in mind we can use Eq. (45) to give an *order of magnitude estimate* of the "upper bound" on $|f|$ as a function of M_N, choosing a value for Λ_C (See Fig. 8). Alternatively, Eq. (45) gives a lower bound on Λ_C as a function of M_N, assuming $|f| = 1$, (see Fig. 9). We can see that the "lower bound" on the compositeness scale coming from $0\nu\beta\beta$ decays is rather weak: $\Lambda_C > 0.3$ TeV at $M_N = 1$ TeV.

In Table 3 we summarize our bounds for sample values of the excited Majorana neutrino mass. In particular, we see that the choice $|f| \approx 1$ is compatible with bounds imposed by experimental limits on neutrinoless double beta decay rates. We remark that, as opposed to the case of bounds coming from the direct search of excited particles, our constraints on Λ_C and $|f|$ *do not depend* on any assumptions regarding the branching ratios for the decays of the heavy particle.

Let us now compare our result in Eq. (45) with that coming from high energy experiments, c.f. Eq. (19). For $m_{\nu^*} = 180$ GeV (the highest accessible mass at the HERA experiments [18,19] with $\Lambda_C = 1$ TeV, $Br(\nu^* \to \nu W) = 0.61$ [37] and $|c_{W\nu^* e}| = |c_{W\nu^* e}|$ one has:

$$|f| < 61. \tag{46}$$

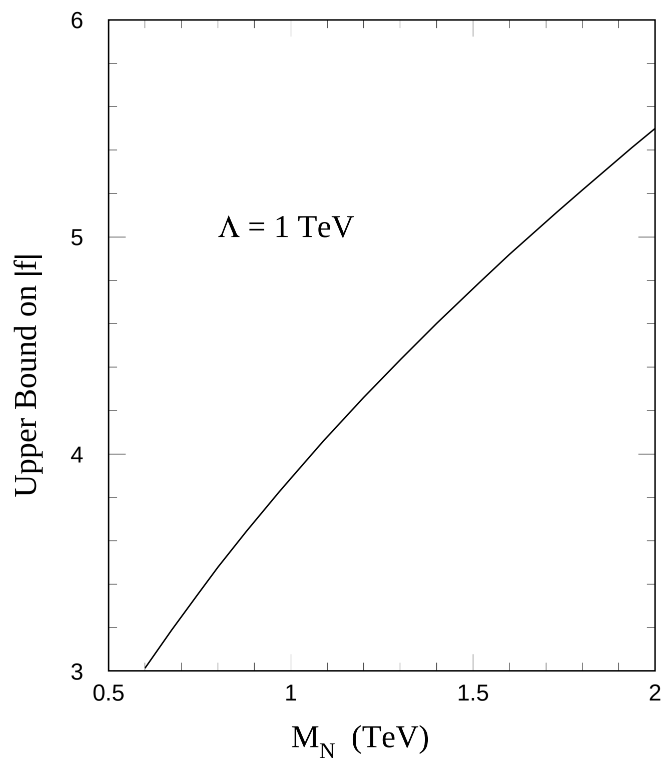

Fig. 8. Most stringent "upper bound" on $|f|$, as it can be derived from Eq. (29), versus the heavy Majorana neutrino mass M_N .

For the same values of $m_{\nu^*} = M_N$ and Λ_C one obtains from the $0\nu\beta\beta$ constraint i.e. Eq. (45):

$$|f| < 1.65 \tag{47}$$

Due to the approximation in the nuclear matrix element discussed above, Eq. (47) represents the most stringent bound that can be derived from $0\nu\beta\beta$ decay and $|f|$ can actually be bigger than 1.65. We can thus conclude that the bounds that can be derived from the low-energy neutrinoless double beta decay are roughly of the same order of magnitude as those coming from the direct search of excited states in high energy experiments.

To obtain more stringent bounds, we need to improve on the measurements of $0\nu\beta\beta$ half-life. However, our bounds c.f. Eq. (40) on ($|f|$ or Λ_C) depend only weakly on the experimental $T_{1/2}$ lower limit ($\propto T_{1/2}^{\pm 1/4}$). To improve the bounds of an order of magnitude we need to push higher, by a factor of 10^4, the lower bound on $T_{1/2}$. We should bear in mind, however, that the simple observation of a few $0\nu\beta\beta$ decay events, while unmistakably proving lepton number violation and the existence of Majorana neutrals, will not be enough to uncover the originating mechanism (including the one discussed here). In order to disentangle the various models, single electron spectra will

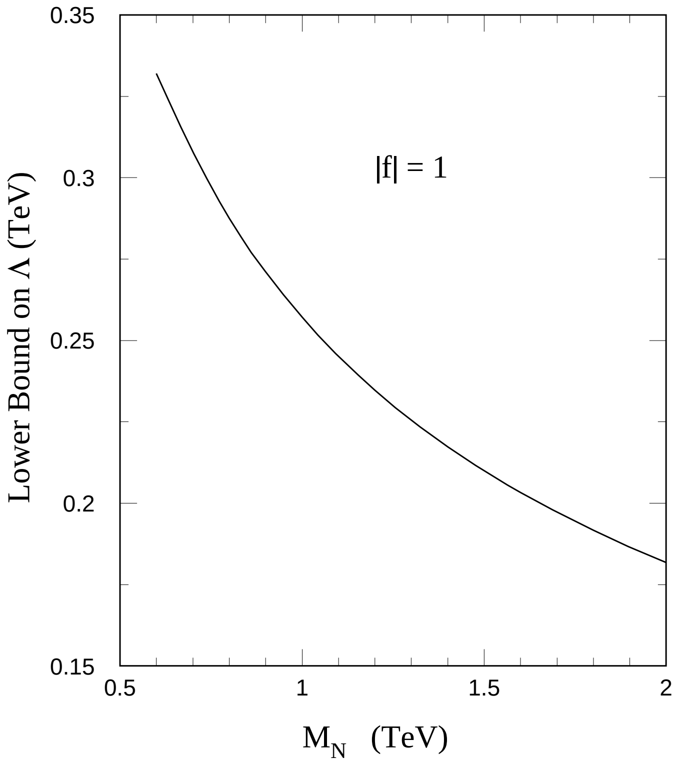

Fig. 9. Most stringent “lower bound” on $\Lambda_{\rm C}$ versus the heavy Majorana neutrino mass M_N, as it can be derived from Eq. (29).

be needed, which would require high statistics experiments and additional theoretical work.

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